

Final-state Polarization in B_s^0 Decays

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Certain $B_s^0 \rightarrow V_1 V_2$ decays (V_i is a vector meson) can be related by flavor SU(3) symmetry to corresponding $B_d^0 \rightarrow V_3 V_4$ decays. In this paper, we show that the final-state polarization can be predicted in the B_s^0 decay, assuming polarization measurements of the B_d^0 decay. This can be done within the scenario of penguin annihilation (PA), which has been suggested as an explanation of the unexpectedly large transverse polarization in $B \rightarrow \phi K^*$. PA is used to estimate the breaking of flavor SU(3) symmetry in pairs of decays. Two of these for which PA makes a reasonably precise prediction of the size of SU(3) breaking are $(B_s^0 \rightarrow \phi\phi, B_d^0 \rightarrow \phi K^{*0})$ and $(B_s^0 \rightarrow \phi \bar{K}^{*0}, B_d^0 \rightarrow \bar{K}^{*0} K^{*0})$. The polarization measurement in the B_d^0 decay can be used to predict the transverse polarization in the B_s^0 decay, and will allow a testing of PA.

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INTRODUCTION

We consider $B \rightarrow V_1 V_2$ decays (V_i is a vector meson). Since the final-state particles are vector mesons, when the spin of these particles is taken into account, this decay is in fact three separate decays, one for each polarization (one longitudinal, two transverse). Naively, within the standard model (SM), the transverse amplitudes are suppressed by a factor of size m_V/m_B (V is one of the vector mesons) with respect to the longitudinal amplitude. Then one expects the fraction of transverse decays, f_T , to be much less than the fraction of longitudinal decays, f_L .

However, it was observed that these two fractions are roughly equal in the decay $B \rightarrow \phi K^*$: $f_T/f_L \simeq 1$ [1, 2, 3]. A similar effect was later seen in $B \rightarrow \rho K^*$ decays [4].

If one goes beyond the naive SM, there are two explanations [5] which account for this “polarization puzzle.” The first is penguin annihilation (PA) [6]. $B \rightarrow \phi K^*$ receives penguin contributions, $\bar{b}\mathcal{O}s\bar{q}\mathcal{O}q$, where $q = u, d$ (\mathcal{O} are Lorentz structures, and color indices are suppressed). With a Fierz transformation, these operators can be written as $\bar{b}\mathcal{O}'q\bar{q}\mathcal{O}'s$. A gluon can now be emitted from one of the quarks in the operators, and can then produce a pair of s, \bar{s} quarks. These then combine with the \bar{s}, q quarks to form the final states ϕK^{*+} ($q = u$) or ϕK^{*0} ($q = d$). These are annihilation contributions. Normally such terms are expected to be small as they are higher order in the $1/m_b$ expansion, and thus ig-

nored. However, within QCD factorization (QCdf) [7], it is plausible that the coefficients of these terms are large [6]. (Within perturbative QCD [8], the penguin annihilation is calculable and can be large, though it is not large enough to explain the polarization data in $B \rightarrow \phi K^*$ [9].)

In QCdf, due to the appearance of endpoint divergences, PA is not calculable, but is modeled [7]. These divergences are regulated with a cut-off, introducing several arbitrary parameters. There is therefore an enormous uncertainty in the size of the PA amplitude as one varies these unknown parameters within certain chosen limits [10].

It is also possible within QCdf that the transverse amplitudes receive significant contributions from perturbative rescattering from charm intermediate states. However, the transverse amplitudes could be purely dominated by PA. In this paper we explore the consequences of the scenario in which PA contributions are large and dominant to see what type of testable predictions result.

The second SM explanation is rescattering [11, 12]. The idea is that nonperturbative rescattering effects involving charm intermediate states, generated by the operator $\bar{b}\mathcal{O}'c\bar{c}\mathcal{O}'s$, can produce large transverse polarization in $B \rightarrow \phi K^*$. A particular realization of this scenario is the following [11]. Consider the decay $B^+ \rightarrow D_s^{*+} \bar{D}^{*0}$. Since the final-state vector mesons are heavy, the transverse polarization can be large. The state $D_s^{*+} \bar{D}^{*0}$ can now rescatter to ϕK^{*+} . If the transverse polarization T is not reduced in the scattering process,

this will lead to $B^+ \rightarrow \phi K^{*+}$ with large f_T/f_L . (A similar rescattering effect can take place for $B_d^0 \rightarrow \phi K^{*0}$.)

It is important to test these explanations in order to determine whether new physics is or is not present. The polarization puzzle has been mainly seen in $\bar{b} \rightarrow \bar{s}$ transitions. However, if PA or rescattering is the true explanation, one also expects to observe large f_T/f_L in $\bar{b} \rightarrow \bar{d}$ decays. In Ref. [5] such decays were discussed, and it was observed that the most promising transitions were those which are dominated by penguin amplitudes. The $\bar{b} \rightarrow \bar{s}$ and $\bar{b} \rightarrow \bar{d}$ penguin decays are

$$\begin{aligned} \bar{b} \rightarrow \bar{s}s\bar{s} \text{ and } \bar{b} \rightarrow \bar{s}d\bar{d} : & B_d^0 \rightarrow \phi K^{0*} \\ & B_s^0 \rightarrow \phi\phi, K^{0*}\bar{K}^{0*} \\ & B^+ \rightarrow \phi K^{*+}, \rho^+ K^{*0} \\ \bar{b} \rightarrow \bar{d}s\bar{s} \text{ and } \bar{b} \rightarrow \bar{d}d\bar{d} : & B_d^0 \rightarrow \bar{K}^{0*} K^{0*} \\ & B_s^0 \rightarrow \phi\bar{K}^{0*} \\ & B^+ \rightarrow K^{*+}\bar{K}^{0*} \end{aligned} \quad (1)$$

(Decays which also receive tree contributions are not included in the current analysis.)

Now, all of these decays are the same under flavor SU(3), which treats d , s and u quarks as equal. The idea is that, given a measurement of the polarization in one decay, one can predict the polarization in another decay using PA or rescattering. However, in relating the two decays, the effect of SU(3) breaking must be included. We can relate the transverse amplitudes of SU(3)-related decays in the scenario in which PA dominates these amplitudes. On the other hand, this relation is unknown in rescattering, which involves long-distance contributions. For this reason, in this paper we consider only PA.

We note that the transverse ($A_{\parallel,\perp}$) and helicity amplitudes (A_{\pm}) are related by $A_{\parallel,\perp} = (A_+ \pm A_-)/\sqrt{2}$. However, A_- for B decays (\bar{A}_+ for \bar{B} decays) has an extra spin-flip $O(1/m_b)$ suppression with respect to A_+ (\bar{A}_-). Consequently, we neglect A_- (\bar{A}_+) and henceforth define $|A_T|^2 \equiv |A_{\parallel}|^2 + |A_{\perp}|^2 \approx |A_+|^2$ (i.e. A_T includes both transverse amplitudes).

First, consider the pair of decays $B_s^0 \rightarrow \phi\phi$ and $B_d^0 \rightarrow \phi K^{0*}$. The main PA contributions to the transverse amplitudes are (the penguin-annihilation term arises only from penguin operators with an internal \bar{t} quark)

$$\begin{aligned} \mathcal{A}_T(B_s^0 \rightarrow \phi\phi) &= |V_{tb}^* V_{ts}| f_{B_s^0} f_{\phi}^2 [2(b_3^{(\phi\phi)} + b_4^{(\phi\phi)})], \\ \mathcal{A}_T(B_d^0 \rightarrow \phi K^{0*}) &= |V_{tb}^* V_{ts}| f_{B_d^0} f_{\phi} f_{K^{0*}} [b_3^{(\phi K^{0*})}], \end{aligned} \quad (2)$$

where $b_3^{(V_1 V_2)}$ and $b_4^{(V_1 V_2)}$ are the QCDf terms corresponding to PA [10]. Throughout the paper, we have dropped the overall factor of $G_F/\sqrt{2}$. (Absolute values are taken for the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements because we are not interested in CP-violating observables here, but rather the rate.)

As noted earlier, QCDf also contains the (perturbative rescattering) term α_4 which can contribute here. It is

the coefficient of the $(V-A) \otimes (V-A)$ piece in the operator product expansion. However, the effect of α_4 could be small in the transverse amplitude because of spin flips. Consequently, it is reasonable to ignore α_4 in pure penguin decays, and we do so here.

The annihilation coefficient b_3 is given by [10]

$$b_3^{(V_1 V_2)} = \frac{C_F}{N_c^2} [(C_5 + N_c C_6) A_3^f + C_5 A_3^i + C_3 A_1^i], \quad (3)$$

where N_c is the number of colors, the C_i are Wilson coefficients, and the incalculable infrared divergences are found in the $A_k^{i,f}$ (due to the endpoint singularity of the final-state distribution amplitudes). The superscripts ‘ i ’ and ‘ f ’ refer to gluon emission from the initial- and final-state quarks, respectively. The subscript ‘1’ refers to the Dirac structure $(V-A) \otimes (V-A)$, while ‘3’ refers to the Dirac structure $(-2)(S-P) \otimes (S+P)$.

The annihilation coefficient b_4 is

$$b_4^{(V_1 V_2)} = \frac{C_F}{N_c^2} (C_4 A_1^i + C_6 A_2^i). \quad (4)$$

Here, $A_2^i \approx A_1^i$ has a suppression factor of $O(1/m_b^2)$ compared to the $A_3^{i,f}$ [10]. Thus, $b_4^{(V_1 V_2)}$ and the third term in $b_3^{(V_1 V_2)}$ can be neglected in Eq. (2).

Now, if we assume that the term containing A_3^f in $b_3^{(V_1 V_2)}$ dominates over the others (as we will see below, the error on this approximation is at the level of only a few percent), we find that

$$\frac{\mathcal{A}_T(B_s^0 \rightarrow \phi\phi)}{2 \mathcal{A}_T(B_d^0 \rightarrow \phi K^{0*})} = \frac{f_{B_s^0} f_{\phi}}{f_{B_d^0} f_{K^{0*}}} \frac{A_3^{f(\phi\phi)}}{A_3^{f(\phi K^{0*})}}, \quad (5)$$

where $A_3^{f(V_1 V_2)}$ has the following integral form [10]:

$$A_3^{f(V_1 V_2)} = \pi \alpha_s \times \int_0^1 dx dy \left\{ \frac{2m_1}{m_2} r_{\chi}^{V_2} \frac{\phi_{a1}(x) \phi_2^{\perp}(y)}{x \bar{y}^2} + \frac{2m_2}{m_1} r_{\chi}^{V_1} \frac{\phi_1^{\perp}(x) \phi_{b2}(y)}{x^2 \bar{y}} \right\}. \quad (6)$$

Here $\phi^{\perp}(\phi_{a,b})$ is the twist-2 (twist-3) light-cone distribution amplitude, x (y) stands for the momentum fraction carried by the quark in V_1 (V_2), and $r_{\chi}^V = (2m_V/m_b) f_V^{\perp}/f_V$. f_V^{\perp} is defined as [13]

$$\langle V(p, \epsilon^*) | \bar{q} \sigma_{\mu\nu} q' | 0 \rangle = f_V^{\perp} (p_{\mu} \epsilon_{\nu}^* - p_{\nu} \epsilon_{\mu}^*). \quad (7)$$

It is useful to make a comment concerning the appropriate interpretation of Eq. (6): if a \bar{K}^* is one of the final particles, the argument of the distribution amplitudes (DAs) is the momentum fraction corresponding to the s quark, whereas for a K^* meson, it is the momentum fraction of the d quark. That is, the DAs are defined as $\phi_{K^*}(z) = \phi_{\bar{K}^*}(\bar{z})$, where $\bar{z} \equiv 1 - z$.

In order to estimate the ratio of the transverse amplitudes in Eq. (5), we need to know the amount of SU(3) breaking in the ratio of $A_3^{f(\phi\phi)}$ and $A_3^{f(\phi K^{0*})}$. To do this, further assumptions are necessary. In what follows, we adopt the same assumptions as those used by the authors of Ref. [10] to carry out their study:

1. the asymptotic form of the light-cone distribution (LCD) amplitudes,
2. a universal parametrization of the end-point singularities, i.e. independent of any particular decay mode,
3. a modeling of the singularities.

It is the first point about which there has been some debate. Certain references have calculated higher-order moments in the LCDs, suggesting that such “non-asymptotic LCDs” are important for some light mesons [14]. If so, then SU(3) breaking in these non-asymptotic pieces will also contribute to the ratio in Eq. (5). Unfortunately, this SU(3) breaking is not calculable, in which case our analysis below will not hold. Equally unfortunately, it is very difficult for experiment to determine which type of LCD is present [15]. Thus, the reader should be aware that our predictions are not only a test of PA dominance, but also of asymptotic LCDs.

Now, the DAs for the final states are universal in the asymptotic limit, i.e. $\phi^\perp(x) = 6x\bar{x}$ and $\phi_a(x) = \phi_b(\bar{x}) = 3\bar{x}^2$. Thus, in this approximation, we find that $A_3^{f(\phi\phi)}$ and $A_3^{f(\phi K^{0*})}$ have exactly the same dependence on x and y [Eq. (6)]. In this case, the ratio between the transverse amplitudes becomes

$$\frac{\mathcal{A}_T(B_s^0 \rightarrow \phi\phi)}{2 \mathcal{A}_T(B_d^0 \rightarrow \phi K^{0*})} = \frac{f_{B_s^0} f_\phi}{f_{B_d^0} f_{K^{0*}}} \frac{2r_\chi^\phi}{\left(\frac{m_{K^{0*}} r_\chi^\phi}{m_\phi} + \frac{m_\phi r_\chi^{K^{0*}}}{m_{K^{0*}}}\right)}, \quad (8)$$

where $f_{B_s^0}/f_{B_d^0} = 1.22 \pm 0.03$ [16], $f_\phi = 221 \pm 3$ MeV [13], and $f_{K^{0*}} = 218 \pm 4$ MeV [13]. The values of f_χ^\perp in r_χ^\perp are theoretically estimated. We have [13] $f_\phi^\perp = f_{K^{0*}}^\perp = 175 \pm 25$ MeV. Since the decay constants and the meson masses are known, the SU(3) breaking in Eq. (8) is well-controlled for this pair.

What this says is that the transverse polarization amplitude in $B_s^0 \rightarrow \phi\phi$ is predicted by PA to be related to that in $B_d^0 \rightarrow \phi K^{0*}$ through Eq. (8). Thus, once one makes the polarization measurement in the B_d^0 decay, one can test PA by making the equivalent measurement in the B_s^0 decay.

The key ingredient in the above analysis is to take two decays in which the final states have the same dependence on the momentum fractions x and y . However, although the pair of decays considered above is the most promising for the analysis, it is not unique. In fact, all decays in a special class have the same dependence on x and y . This class contains $\bar{b}(B_d^0, B^+) \rightarrow \bar{s}d\bar{d}$ (the parenthesis indicates that this transition includes only B_d^0 or B^+ , and not B_s^0 , decays) and $\bar{b} \rightarrow \bar{s}s\bar{s}$ penguin decays. In addition, only decays to ground-state spin-1 mesons are included (excited mesons have different DAs in general). Thus, excluding those decays which also receive tree contributions, these correspond to $B^+ \rightarrow \rho^+ K^{0*}$, $B_d^0 \rightarrow \phi K^{0*}$,

$B^+ \rightarrow \phi K^{+*}$, and $B_s^0 \rightarrow \phi\phi$. The important point here is that the dependences on the momentum fractions in A_3^f are the same for **every** decay belonging to this class. Therefore, in the comparison of any two of these decays, the integrals containing the singularities cancel in the ratio of the transverse amplitudes, and this even before using a cutoff to regulate the end-point divergences.

Of course, the size of the SU(3) breaking will depend on the pairs of decays considered [see Eq. (8)]. For example, we expect the pair $B_s^0 \rightarrow \phi\phi$ and $B^+ \rightarrow \phi K^{+*}$ to be as good as $B_s^0 \rightarrow \phi\phi$ and $B_d^0 \rightarrow \phi K^{0*}$. However, the SU(3) breaking in $B^+ \rightarrow \rho^+ K^{0*}$ and $B^+ \rightarrow \phi K^{+*}$ also turns out to be about the same size. In all cases, the SU(3) breaking can be worked out as we have done above.

Another class in which the final states have the same dependence on x and y is given by the decays $\bar{b}(B_s^0) \rightarrow \bar{d}d\bar{d}$ and $\bar{b}(B_s^0) \rightarrow \bar{s}d\bar{d}$. Even if one of the most promising B_s^0 decay modes to be measured in the near future, $B_s^0 \rightarrow K^{0*}\bar{K}^{0*}$, belongs to this class, we cannot use it to make predictions because its decay-class partner, $B_s^0 \rightarrow \rho^0 \bar{K}^{*0}$, typically receives tree contributions. Therefore, this class is not very useful since it only contains one pure penguin decay.

The third class is defined by the transition $\bar{b} \rightarrow \bar{d}s\bar{s}$. This includes the decays $B_s^0 \rightarrow \phi \bar{K}^{0*}$ and $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$, and this pair is particularly promising. The prediction for the ratio of the transverse amplitudes is

$$\frac{\mathcal{A}_T(B_d^0 \rightarrow \bar{K}^{0*} K^{0*})}{\mathcal{A}_T(B_s^0 \rightarrow \phi \bar{K}^{0*})} = \frac{f_{B_d^0} f_{K^{0*}}}{f_{B_s^0} f_\phi} \frac{2r_\chi^{K^{0*}}}{\left(\frac{m_{K^{0*}} r_\chi^\phi}{m_\phi} + \frac{m_\phi r_\chi^{K^{0*}}}{m_{K^{0*}}}\right)}. \quad (9)$$

The pair $B_s^0 \rightarrow \phi \bar{K}^{0*}$ and $B^+ \rightarrow K^{+*} K^{0*}$ can be treated similarly.

An important consequence of the above discussion is that it is not possible to get cancellations in the A_3^f ratios through the comparison of decays belonging to different classes. For example, we will not obtain a clean result by comparing $B_s^0 \rightarrow K^{0*} \bar{K}^{0*}$ and $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$, even though the final states are identical in the two decays. In addition, we have presented the list of decays in Eq. (1) in terms of $\bar{b} \rightarrow \bar{d}$ or $\bar{b} \rightarrow \bar{s}$ transitions, motivated by the apparent predominant role played by the $\bar{b} \rightarrow \bar{s}$ transitions in the polarization puzzle. However, as shown above, this is not the most natural way to classify the decays in order to achieve the most reliable predictions for penguin annihilation within QCDf.

We now turn to the estimation of errors in Eqs. (8) and (9) due to the single inclusion of A_3^f in the transverse amplitudes. We study the relative magnitude of the $A_{1,3}^i$ terms in b_3 [Eq. (3)], as well as the relevance of the annihilation coefficient b_4 [Eq. (4)] in the cases it is appropriate. To do this, we follow several applications of QCDf in which the incalculable infrared divergences in the $A_k^{i,f}$ are isolated with a cutoff and the penguin annihilation amplitude is modeled by introducing unknown parameters.

In passing, we note the following. Previously, we mentioned that we use the same assumptions as those in Ref. [10]. Although we work within the same restricted theoretical framework as this reference, and although it is true that the size of the error in our predictions could be affected by large uncertainties related to the choice of this particular scenario, we emphasize that our results go beyond the analysis made in Ref. [10]. There, due to the parametrization of the infinities, the uncertainties in the individual transverse amplitudes turn out to be at the level of one hundred percent for most of the decay channels [19]. Instead, we show here that it is possible to obtain more accurate predictions with the same theoretical inputs used in the treatment of the divergent integrals when specific decays are compared.

To illustrate the procedure used in the estimation of errors, we focus on the first pair, $B_s^0 \rightarrow \phi\phi$ and $B_d^0 \rightarrow \phi K^{0*}$ (a similar reasoning holds for $B_s^0 \rightarrow \phi \bar{K}^{0*}$ and $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$, as well as several other decay pairs). We first write the $A_{1,3}^{f,i}$ after the parametrization of the infrared divergences is applied [10]

$$\begin{aligned} A_{1,2}^i &\approx 18\pi\alpha_s \frac{m_1 m_2}{m_B^2} \left(\frac{1}{2} X_L + \frac{5}{2} - \frac{\pi^2}{3} \right), \\ A_3^i &\approx 18\pi\alpha_s \left(\frac{m_1}{m_2} r_\chi^{v_2} - \frac{m_2}{m_1} r_\chi^{v_1} \right) (X_A^2 - 2X_A + 2), \\ A_3^f &\approx 18\pi\alpha_s \left(\frac{m_1}{m_2} r_\chi^{v_2} + \frac{m_2}{m_1} r_\chi^{v_1} \right) (2X_A^2 - 5X_A + 3), \end{aligned} \quad (10)$$

We see that $A_{1,2}^i$ and A_3^f are symmetric in the interchange of $V_1 \leftrightarrow V_2$, while A_3^i is antisymmetric. X_A and X_L contain the same input parameters, but have different end-point-divergence behavior:

$$\begin{aligned} X_A &= (1 + \rho_A e^{i\phi_A}) \ln \frac{m_B}{\Lambda_h}, \\ X_L &= (1 + \rho_A e^{i\phi_A}) \frac{m_B}{\Lambda_h}. \end{aligned} \quad (11)$$

Here, Λ_h is an input parameter ($\Lambda_h = 0.5$ GeV [10]), and ϕ_A is an arbitrary phase. We have taken Λ_h , ϕ_A , and ρ_A to be the same for every decay mode.

To study the relative significance of b_4 and the neglected terms in b_3 , we evaluate the following ratios [Eqs. (3) and (4)]:

$$\begin{aligned} r_{b_4}^{(V_1 V_2)} &= \frac{C_4 + C_6}{C_5 + N_c C_6} \frac{A_{1,2}^{i(V_1 V_2)}}{A_3^{f(V_1 V_2)}}, \\ r_{b_3}^{(V_1 V_2)} &= \frac{C_3}{C_5 + N_c C_6} \frac{A_1^{i(V_1 V_2)}}{A_3^{f(V_1 V_2)}}, \\ R_{b_3}^{(\phi K^{0*})} &= \frac{C_5}{C_5 + N_c C_6} \frac{A_3^{i(\phi K^{0*})}}{A_3^{f(\phi K^{0*})}}, \end{aligned} \quad (12)$$

where $V_1 V_2 = \phi\phi$, ϕK^{0*} (note that $R_{b_3}^{(\phi\phi)}$ is zero).

Although the values of these ratios are quite uncertain, in large part due to the (arbitrary) value of Λ_h , in virtually all cases it is found that $|r_{b_4}^{(V_1 V_2)}| \lesssim \mathcal{O}(10^{-2})$ and

$|r_{b_3}^{(V_1 V_2)}|, |R_{b_3}^{(\phi K^{0*})}| \lesssim \mathcal{O}(10^{-3})$. One can see this as follows. First, the ratios of the relevant Wilson coefficients in Eq. (12) at $\mu = m_b/2$ are as follows [7]:

$$\begin{aligned} \frac{C_4 + C_6}{C_5 + N_c C_6} &\approx 0.63, \\ \frac{C_3}{C_5 + N_c C_6} &\approx -0.11, \quad \frac{C_5}{C_5 + N_c C_6} \approx -0.05, \end{aligned} \quad (13)$$

Second, we have $(A_1^{i(V_1 V_2)} / A_3^{f(V_1 V_2)}) \sim \mathcal{O}(m_1 m_2 / m_B^2)$, and $|A_3^{i(\phi K^{0*})} / A_3^{f(\phi K^{0*})}| \sim |(a - b) / (a + b)| \approx 0.07$ (where a denotes $(m_{K^{0*}} / m_\phi) r_\chi^\phi$ and b is given by $K^{0*} \leftrightarrow \phi$).

We evaluate the three ratios in Eq. (12) by considering many different values in the ranges $0 \leq \rho_A \leq 2$ and $0 \leq \phi_A \leq 2\pi$. We find that $|r_{b_4}^{(V_1 V_2)}|, |r_{b_3}^{(V_1 V_2)}|, |R_{b_3}^{(\phi K^{0*})}| \ll 1$ always, except for a singular behavior at $\phi_A = 0, 2\pi$. The largest contribution to the error arises from $r_{b_4}^{(V_1 V_2)}$ but it remains at the level of a few percent within the scanned region of the parameter space. Thus, we have covered a wide set of models of the infrared singularities. The point here is that, although the precise values of the ratios are very uncertain, they are always small.

We therefore conclude that the PA dominance hypothesis leads to a clean prediction for the ratio of transverse amplitudes in the pair $B_s^0 \rightarrow \phi\phi$ and $B_d^0 \rightarrow \phi K^{0*}$ [Eq. (8)], though this result is a direct consequence of the particular modeling of the suppressed terms (e.g. asymptotic LCDs). We have also analyzed the pair $B_s^0 \rightarrow \phi \bar{K}^{0*}$ and $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$, as well as the pairs containing the charged modes and $(B^+ \rightarrow \rho^+ K^{0*}, B^+ \rightarrow \phi K^{*+})$, under the same set of assumptions, and we have obtained similar conclusions.

We have therefore seen that there are a number of decay pairs within a given class whose transverse polarizations are related. Some of these pairs involve a B_s^0 decay. Now, the B -factories BaBar and Belle have made many measurements of B_d^0 and B^+ mesons. But it is only relatively recently, at hadron colliders, that B_s^0 mesons have started to be studied. This will increase when the LHCb turns on. It should be possible to make measurements of the transverse polarization in some B_s^0 decays in the near future, and to test the PA/QCDf hypothesis.

Above, we have presented the ratio of \mathcal{A}_T 's for two pairs of decays. However, it is perhaps better to present a ratio of f_T 's since this is what will actually be measured. \mathcal{A}_T and f_T are related by including information about the branching ratio (BR): $f_T = |\mathcal{A}_T|^2 / (\Gamma BR/PS)$, where Γ

is the total width and PS is the phase space. We find

$$\begin{aligned} \frac{f_T(B_s^0 \rightarrow \phi\phi)}{f_T(B_d^0 \rightarrow \phi K^{0*})} &= \\ 3.22 \pm 0.72 \frac{BR(B_d^0 \rightarrow \phi K^{0*})}{BR(B_s^0 \rightarrow \phi\phi)}, \\ \frac{f_T(B_d^0 \rightarrow \bar{K}^{0*} K^{0*})}{f_T(B_s^0 \rightarrow \phi \bar{K}^{0*})} &= \\ 0.62 \pm 0.12 \frac{BR(B_s^0 \rightarrow \phi \bar{K}^{0*})}{BR(B_d^0 \rightarrow \bar{K}^{0*} K^{0*})}. \end{aligned} \quad (14)$$

The numbers have been obtained by taking values of masses and lifetimes from the Particle Data Group without errors [17], along with the theoretical estimates of the decay constants given above. The predictions given in Eq. (14) will yield a test of PA. If there are discrepancies in the measurements, this may indicate the presence of new physics, in $\bar{b} \rightarrow \bar{s}$ and/or $\bar{b} \rightarrow \bar{d}$ transitions.

The decays $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$ [18] and $B_d^0 \rightarrow \phi K^{0*}$ [2] have both been measured, so that this information can be included in Eq. (14):

$$\begin{aligned} B_d^0 \rightarrow \bar{K}^{0*} K^{0*} : f_L &= 0.81_{-0.12}^{+0.10} \pm 0.06, \\ BR &= (0.96_{-0.19}^{+0.21}) \times 10^{-6}, \\ B_d^0 \rightarrow \phi K^{0*} : f_L &= 0.49 \pm 0.04, \\ BR &= (9.5 \pm 0.8) \times 10^{-6}. \end{aligned} \quad (15)$$

f_T is defined as $f_T = 1 - f_L$.

To summarize, a large f_T/f_L has been observed in several $\bar{b} \rightarrow \bar{s}$ $B \rightarrow V_1 V_2$ decays. There are two explanations of this measurement within the standard model – penguin annihilation (PA) and rescattering. Now, one logically also expects to see a large f_T/f_L in certain $\bar{b} \rightarrow \bar{d}$ decays. The most promising decays are those dominated by penguin amplitudes, and there are quite a few $\bar{b} \rightarrow \bar{d}$ and $\bar{b} \rightarrow \bar{s}$ penguin decays. All of these are equal under flavor SU(3) symmetry. Given the measurement of f_T/f_L in one decay, if one wishes to predict f_T/f_L in another decay, it is necessary to take SU(3) breaking into account. However, it is only within a specific scenario of PA dominance for the transverse amplitudes, and for special classes of decay pairs, that this SU(3) breaking can be estimated. We therefore assume that it is PA alone which is the source of the large transverse polarization and explore its consequences.

We find that there are several decay pairs for which PA makes a reasonably precise estimate of the SU(3) breaking (assuming asymptotic LCDs). Thus, given the measurement of f_T/f_L in one decay, PA makes a prediction for the transverse polarization in the second decay. In this paper we have concentrated on two decay pairs that involve B_s^0 mesons: ($B_s^0 \rightarrow \phi\phi$, $B_d^0 \rightarrow \phi K^{0*}$) and ($B_s^0 \rightarrow \phi \bar{K}^{0*}$, $B_d^0 \rightarrow \bar{K}^{0*} K^{0*}$). The polarization measurement in the B_d^0 decay allows one to predict the transverse polarization in the B_s^0 decay. This will permit

the explicit testing of PA, probably in the near future at the LHCb.

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